Polymer Chains and Baryons in a Strongly Coupled Quark-Gluon Plasma

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Recently there was a significant change of views on physical properties and underlying dynamics of Quark-Gluon Plasma at $T = 170-350\,MeV$, produced in heavy ion collisions at RHIC. Instead of being a gas of q, g quasiparticles, a near-perfect liquid is observed. Also, precisely in this temperature interval, the interaction deduced from lattice studies is strong enough to support multiple binary bound states. This work is the first variational study of *multibody* bound states. We will consider: (i) "polymer chains" of the type $\bar{q}gg..gq$; (ii) baryons (qqq); (iii) closed (3-)chains of gluons (ggg). We found that chains (i) form in exactly the same T range as binary states, with the same binding *per bond*. The binding and T-range for diquarks, baryons and closed 3-chains are also established. We point out that the presence of chains, or possibly even a chain network, may drastically change the transport properties of matter, such as charm diffusion or jet energy loss. We further suggest that it seems to exist only for $T = (1-1.5)T_c$ and thus there may be a "latent period" for charm/jet quenching in RHIC collisions, while matter cools down to such T.

1. Introduction

Quark-gluon plasma (QGP) is a high-temperature phase of QCD, and the word "plasma" means the same sense as in electrodynamic plasma physics, the presence of free (color) charges. QGP is different from hadronic phase because charges are screened by the medium [1] rather than confined in neutral objects. Lattice simulations have shown, that QGP exists above the phase transition $T > T_c \approx 170\,MeV$, and is not only deconfined but also possesses the restored chiral symmetry.

Asymptotic freedom of non-Abelian gauge theories like QCD ensures that for high enough temperature $T >> \Lambda_{QCD}$ this phase becomes *weakly* coupled (wQGP), with most of interactions characterized by small coupling $\alpha_s(p \sim T) << 1$. In this domain wQGP is essentially a near-ideal gas of its fundamental constituents, quarks and gluons.

QGP is experimentally studied via heavy ion collisions, at CERN SPS and last years at BNL RHIC collider, reaching temperatures up to about $T \approx 2T_c$. Success of hydrodynamical description [2] of observed collective flows has indicated, that all dissipative lengths are very short and thus the produced matter cannot be a weakly coupled gas but rather a very good liquid [3]. Recent studies of charm transport [4] and preliminary RHIC data indicate, that charm diffusion constant is also much smaller than pQCD predictions. This complements another well known puzzle, of unexpectedly strong jet quenching at RHIC,

with its so far unexplained angular dependence.

In order to explain all of these features, a radically new picture of QGP at such temperatures is being developed, referred to as *strongly coupled* QGP, or sQGP. It has been in particular pointed out in [5] that the interaction is strong enough to preserve the meson-like bound states up to about $T = 2T_c \sim 300\, MeV$, (the temperature range corresponding to QGP at RHIC), although in a strongly modified form. It was then pointed out in [6] that also multiple binary *colored* bound states should exist in the same T domain. Since QGP is a deconfined phase, there is nothing wrong with that, and the forces between say singlet $\bar{q}q$ and octet qg quasiparticle pairs are about the same. (For potential-like forces the Casimir scaling gives 9/8 ratio, for string-like ones the ratio is just 1.) Some of those states (charmonium) were observed on the lattice [7] at T up to about $2.5T_c$, while existence of most of these states, especially colored, still has to be checked.

By this work we make the first step toward the understanding of the *multibody* bound states. For definiteness, we will use similar parameterized lattice-based interactions as in [6].

In the QCD vacuum, the potential for two color charges is traditionally written as a sum of a Coulomb and linear potential, dominating small and large distances, respectively. At $T > T_c$, by definition of the deconfined phase, the effective string tension vanishes and the potentials go to a constant at $r \to \infty$. But that does not mean that string-like field configurations of color field disappear right at T_c: as explained by Polyakov long ago [8], the string tension which vanishes at $T = T_c$ should be understood as the free energy, $F = V - TS$, while the string energy V and its entropy (related to the number of configurations) S are finite but cancelling each other. This picture of the deconfinement suggests by itself that the "mixed phase", at $T = T_c$, contains a lot of very long strings. It is natural to think then, that strong (although finite-range) interaction between the charges at $T > T_c$ is also related with strings.

Applications of lattice-based binary potentials for static quarks to multi-body problems meets a fundamental question: what part of this interaction is (i) of a "potential-type" or (ii) of a "string-type". In the former case the potential energy of a manybody state is the sum over *all* pairs of particles, while in the latter only some pre-defined partners, "connected by a string" are allowed to interact. The issue is well known and was discussed in literature for baryons for decades. In Fig.1(left) we show how two pictures look like, with strings in (a) ending at a string junction.

The discussion above motivates us to consider interaction to be string-like in this sense. An additional reason for that is the especially simple multi-particle states appear, namely the *polymeric chains* made of repeated gluons with $\bar{q}$ and q at the ends, see Fig.1(c). (Recall that gluons have 2 color indices and can be viewed as 2 different color charges connected to two strings.)

For baryons, it turns out that the string-based picture (a) and potential one (b) (with the Casimir factor 1/2 compared with meson potential) give very close results, as is the case in vacuum [9]. Also, this is supported by recent lattice study of free energy of static three quark systems[10] which found that the connected part of qqq-singlet free energies above T_c are decomposable into three qq-triplet (diquark) free energies for all distances. We will not discuss more complicated possible structures, like "polymerized" baryons with extra gluons or a network of chains, connected by color junctions.

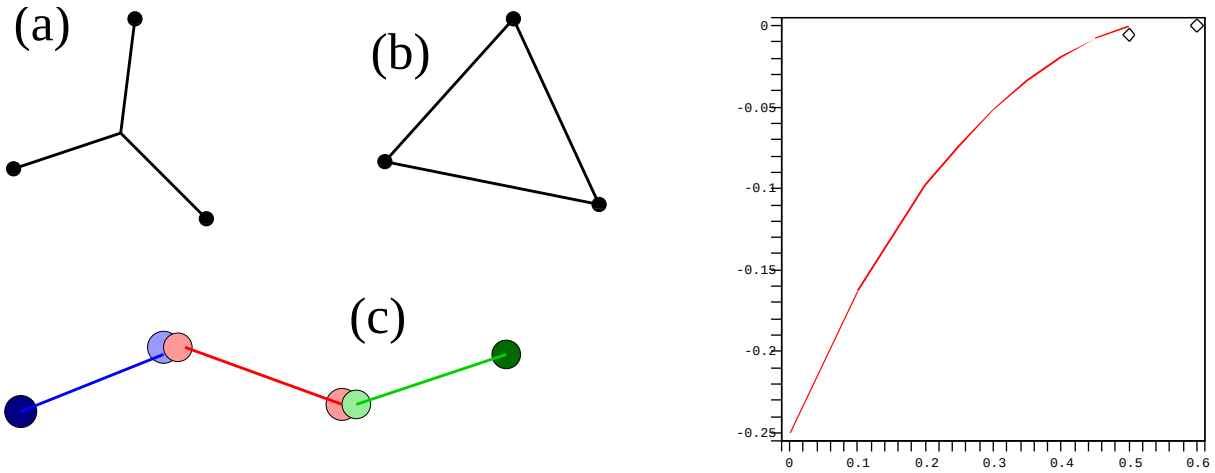

Figure 1. (left) The interaction in baryons for "string-like" interaction (a) versus the "potential-like" interaction (b). The double circles with different colors (online) in (c) represent gluons, and it is an example of 4-chain $\bar{q}ggq$. (right) Dependence of the binding energy on the Debye screening mass for the simple exponential trial functions. The units are explained in the text. Two diamond points indicate positions of the exact solutions.

Although we would not calculate the statistical and transport properties of sQGP as such in this work, we would comment on them at the end in the last section.

2. The coordinates and the variational procedure

We will denote particle coordinates by $\vec{x}_i, i = 1, n$ where the vector means a 3-d space and the index is for the particle number. Since the center of mass coordinate does not appear in the potential energy, the corresponding momentum is conserved and can be put to zero. It is a standard procedure to use some redefined coordinates which reduce the n-body problem to the (n-1) one: those coordinates would be denoted by $\vec{r}_i, i = 1, n$. Other notations we will use are

$$\vec{\partial}_i = \frac{\vec{\partial}}{\partial x_i}, \vec{D}_i = \frac{\vec{\partial}}{\partial r_i}, \tag{1}$$

related by $\vec{\partial}_i = \vec{D}_j M_{ji}$ where M is the coordinate matrix $r = Mx$ (namely $\vec{r}_j = M_{ji}\vec{x}_i$).

For reference and further comparison, let us give explicit form of the kinetic energy for (3-body) Jacobi coordinates

$$\vec{r}_1 = \vec{x}_1 - \vec{x}_2; \vec{r}_2 = (-1/2)\vec{x}_1 + (-1/2)\vec{x}_2 + \vec{x}_3; \\ \vec{r}_3 = (1/3)(\vec{x}_1 + \vec{x}_2 + \vec{x}_3) \tag{2}$$

$$\frac{1}{2M}(\vec{\partial}_i)^2|_{3Jacobi} = \frac{1}{2M}(2\vec{D}_1^2 + (3/2)\vec{D}_2^2 + (1/3)\vec{D}_3^2) \tag{3}$$

Note that it is a diagonal form avoiding coordinate mixing.

However for the polymer problem we found that more useful coordinates are "chain coordinates" defined by the following set of n coordinates

$$\vec{r}_1 = \vec{x}_2 - \vec{x}_1; ...\vec{r}_{n-1} = \vec{x}_n - \vec{x}_{n-1}; \qquad (4)$$
$$\vec{r}_n = \frac{1}{n}(\vec{x}_1 + ... + \vec{x}_n)$$

The corresponding kinetic energies for N-chains is

$$\frac{1}{2M}(\vec{\partial}_i)^2|_{N-chain} = \frac{1}{2M}(2\vec{D}_1^2 + 2\vec{D}_2^2 + \cdots + 2\vec{D}_{N-1}^2 + \qquad (5)$$
$$(1/N)\vec{D}_N^2 - 2\sum_{i=1}^{N-2}\vec{D}_i\vec{D}_{i+1})$$

Although it is not diagonal, it is very simple instead, with the same "reduced mass" in each diagonal term[1] . Furthermore, in "string-like" approach the potential energy is very simple, just a sum over all "bonds" along the chain

$$E_{pot} = V(r_1) + ... + V(r_{n-1}) \qquad (6)$$

Note that we dropped vector notation here: it means that only the lengths of these coordinates matter. Furthermore, all angular variables would not be important for the ground s-wave states to be discussed. If so, the wave function factorizes

$$\Psi(\vec{r}_1...\vec{r}_n) = \prod_{i=1}^{n-1}\psi(r_i) \qquad (7)$$

and as a result the average value of the non-diagonal terms in the kinetic energy would vanish

$$< \vec{D}_i\vec{D}_j > |_{i\neq j} = 0 \qquad (8)$$

Finally, as the diagonal terms have the same reduced mass as a 2-body problem, the problem obviously splits into (n-1) Schroedinger equations. This completes the proof that there is the same binding energy *per bond* (not per particle) as for mesonic states.. The binding *per particle* of course grows, doubling the binding in binary states as the length of the chain grows[2].

3. Mesons and polymers in a variational approach

Since the 2-body problem and (as shown above) polymer chains can be easily solved numerically for any potential, and for relevant lattice-based potential it was already done in [6], the reader may be surprised why we discuss it here. We however found it instructive

[1]This of course requires that all involved quasiparticles, the quark and the gluon ones, have the same mass, which is however approximately fulfilled by available lattice data.

[2]We again remind the reader that it only happened because in string-like approach one can ignore all the interactions between non-next-neighbors along the chains.

to start with a simple example, for which all calculations and integrals are simple and can be done analytically

Let the potential be just a screened Coulomb (or Yukawa) potential

$$V \;=\; -\frac{\alpha\, e^{-M_D\, r}}{r} \tag{9}$$

and the trial function be as simple as possible, namely an exponential function

$$\psi \;=\; e^{-Ar} \tag{10}$$

The average energy is

$$\begin{aligned} < H > \;&=\; < (-1/2/M) * D^2 - (1/r/M) * D + V(r) > \\ &=\; 1/2\, \frac{A^2}{M} - \frac{4 * \alpha * A^3}{(M_D + 2A)^2} \end{aligned} \tag{11}$$

and it can be easily minimized in respect to parameter A. The results (for $\alpha = 1, M = 1/2$) as a function of M_D are plotted in Fig.1(right).

Note that the simple exponential trial function we use is exact for a Coulomb problem ($M_D = 0$ case), but of course is not so for a screened potential. Although the energy may seem to be quite close to exact ones, obtained from numerical solution of the Schroedinger equation, true wave function is not particularly well reproduced by it as M_D grows. In particular, the curve in Fig.1(right) crosses zero at its endpoint in a wrong manner: in fact the curve must have a horizontal tangent at the endpoint. Moreover, the critical value for the level disappearance predicted by exponential trial function

$$\frac{\alpha M}{M_D}|_{zerobinding} = 1 \tag{12}$$

is not at all accurate, as a comparison to exact behavior (indicated by two points in the right upper corner of Fig.1)(right) changes the r.h.s. of (12) to a smaller value, 5/6.

One may ask whether it is possible to test the quality of the trial wave function without a knowledge of the exact result (as would be the case in multi-dimensional problems to be addressed in subsequent section). A well known observable for that is the so called "energy dispersion" variable, defined by

$$d =< H^2 > - < H >^2 \tag{13}$$

If the trial wave function is an eigenvector, d=0, otherwise it characterizes the quality of the approximation. In Fig.2 we show how it depends on M_D. As expected, it is zero for pure Coulomb problem ($M_D = 0$) but strongly grows with M_D, indicating loss of quality of the approximation. Nevertheless, we emphasize that while this energy dispersion is quite sensitive to how close the trial function is to true solution, the energy itself is much less sensitive to the details of the wavefunction shape.

Now we turn to a realistic variational approach for the two-body bonds in mesons and polymers, using a parameterized temperature-dependent potential extracted from lattice data [6]

$$V(T,r) = \;-\; (\frac{4}{3r} + \frac{8T}{3})\frac{e^{-2Tr}}{log(1+3T)}\frac{4T}{r(1+3T)}\frac{e^{-2Tr}}{(log(1+3T))^2} \tag{14}$$

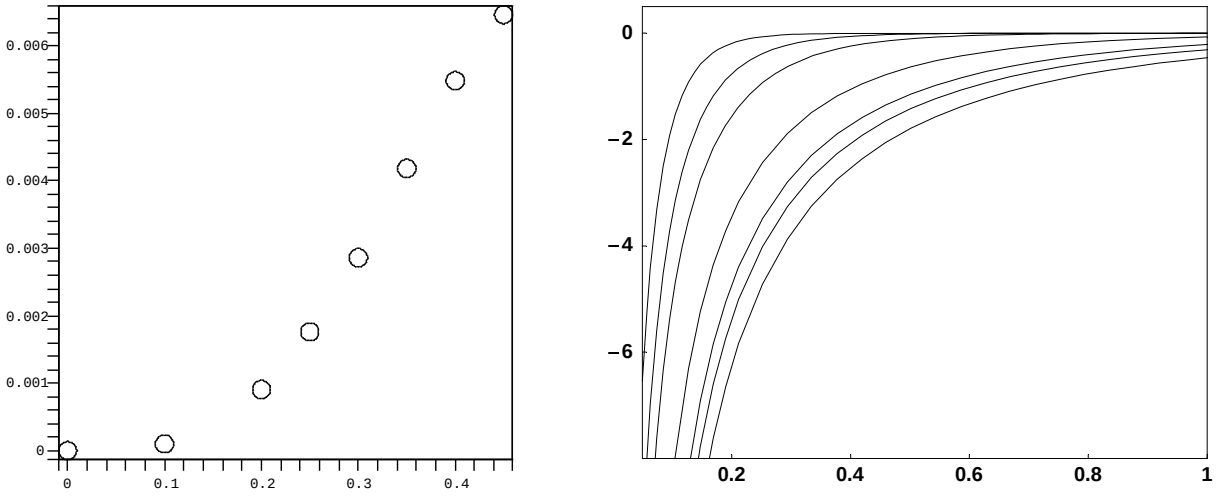

Figure 2. (left) Dependence of the "energy dispersion" variable d on the Debye screening mass M_D, for the simple exponential trial functions. (right) The static potentials $V(T,r)$ (in unit T_c) as a function of the distance r (in unit $1/T_c$). The values of temperature used are $T = 1, 1.2, 1.4, 2, 4, 6, 10\, T_c$ for curves from right to left.

Note that we have scaled all dimensional quantities with proper powers of T_c, namely T gives T/T_c, r means rT_c, and so on. We will keep these units throughout present paper. Plots of the potential at different temperatures are shown in Fig.2(right). At very short distance this potential goes as a Coulomb (as expected from short-range one-gluon exchange interaction) while at very large distance it is a screened Coulomb potential which damps so fast as if it is almost vanishing. The log term is simplified compared to the original form in [6] where it was $log(1/r + 3T)$, the parameterization consistent with the asymptotic freedom at small r.

Note that this is basically for color singlet $\bar{q}-q$ (and for color octet $g-g$ approximately), so we need to add appropriate overall coefficients for other channels like diquark.

According to the features of the above potential, we employ a trial wave function as following (with $\vec{r}$ the 2-body relative coordinate)

$$\phi(r) = e^{-Ar - \frac{C^2 r^2}{(C^2 r^2+1)}(B-A)r - \frac{1}{2} log(C^2 r^2+1)} \tag{15}$$

It has the asymptotic forms to be

$$r \to 0 \ : \ \phi \to e^{-A\, r} r \to \infty \ : \ \phi \to \frac{e^{-B\, r}}{C\, r} \tag{16}$$

Here the parameter C controls the interpolation between the short distance Coulomb behavior and the long distance free particle solution. All the three parameters have the same unit for which we use T_c.

With the (temperature-dependent) potential and trial function at hand, we then find the binding energy at different temperature by minimizing energy (kinetic plus potential

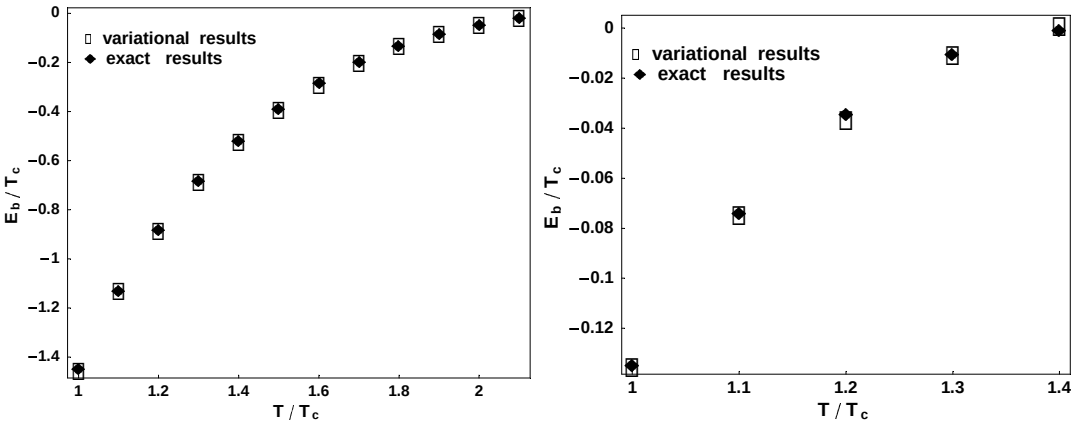

Figure 3. Dependence of the meson (left) and diquark (right) binding energy on the temperature. The units are in T_c.

energy) according to the three variational parameters A, B, C. A non-relativistic form for kinetic part is used, since the color-charged quasiparticles, both quarks and gluons, are found by lattice simulations to be rather heavy at temperatures not so high from T_c, for which we use the same constant value $m^* = 800MeV$ in our calculation for simplicity. We employed the well-known Metropolis Monte Carlo[11] to evaluate the energy. The results are shown in Fig.3. We also plot the exact energies obtained by numerically solving the Schroedinger equations for comparison. As can be seen, the variational results coincide perfectly with the exact results. The energy dispersions are of order 0.01 (not vanishingly small), but as we emphasized before and shown here, the energy itself could still be very accurately evaluated. For the optimal values of wavefunction parameters, A decreases from 3.125 to 2.25 and B from 3.125 to 0.5 with increasing temperature, and C keeps between 0.5-1 to which the energies are less sensitive.

The binding energy E_b is greater than temperature for 1-1.1 T_c and comparable with temperature up to 1.6 $T_c(|E_b|/T$ about 0.26, and $e^{0.26} = 1.3)$. Naively following a $e^{-E_b/T}$ (at T_c as high as 4.3) factor argument we may say the formation of meson-like bonds in this T-region is quite favored. Now recalling the proof in previous section, a polymer chain with N elements will have a binding energy $(N-1)E_b$ (E_b per bond), and further more, the long chains have much more statistical degeneracies (according to vast options arranging intermediate gluons' quantum numbers), so we expect abundance of polymer chains at temperatures just above T_c. And this, as we will suggest, may dramatically contribute to jet quenching as well as transport properties.

4. Baryons at $T > Tc$

After consideration of meson-like structures, it is of course natural and interesting to address the question of possible baryonic bound states above T_c. While lattice calculation found mesonic bound states above T_c[7], there is no available information about baryon up till now. Also it is very important to see the role of baryons in the deconfined phase

of QCD. For example, baryons carry conserved quantum numbers like B, S which may be combined to give useful experimental/lattice signal[16] and contribute more to study of thermodynamics at non-zero chemical potential[12]. So in this section we conduct a variational approach for baryons in similar way as was used in the past to study baryons in vacuum[9].

We consider baryons(anti-baryons) as closed 3-chains of quarks(anti-quarks), which contains 3 pairs of diquarks(anti-diquarks). Thus to study baryons, we first start with diquarks at $T \geq T_c$. Diquarks in the deconfined phase of QCD are of their own importance also. For the diquark channel, the mutual interaction coupling should be one half the quark-anti-quark channel, thus we adjust the potential (14) by a coefficient $C/C_{\bar{q}q} = 1/2$ to use it for diquark bound states. We use the same trial function (15) as in meson case, and then minimize energy according to A, B, C. In Fig.3(right) we plot the diquark binding energy as a function of temperature. The exact numeric results are also presented to compare and justify our variational approach. Diquark states are much more shallowly bound than meson and thus more easily melted.

For baryons, we construct the trial wavefunction on the basis of the three diquark pairs. Since we're only interested in the ground state, it is reasonable to use a totally symmetric s-wave spatial configuration. The color wavefunction should still be the singlet (then antisymmetric among the 3 quarks) to guarantee the attractive interaction, but for spin and flavor part, the only constraint is to be symmetric and there are a lot of ways to arrange it (increasing statistical degeneracies). Particularly, we want emphasize that very different from constructing baryon in vacuum, now s quark is more or less the same as u, d quarks, since its current mass, of the order of T_c is much less than quasiparticle mass and the current mass splitting is now unimportant. So we write down the following

$$\psi(\vec{x}_1, \vec{x}_2, \vec{x}_3) = \phi(r_{12})\phi(r_{23})\phi(r_{31}) \tag{17}$$

Here $\phi(r)$ is from (15) which has been used for meson and diquark states, and $\vec{r}_{ij} = \vec{x}_i - \vec{x}_j$. Again we minimize energy according to the variational parameters A, B, C. Note now for this 3-body object, we have two kinetic energy terms according to the first two in (3) (one term with reduced mass factor 1/2 and the other 2/3), while for potential energy we need count for each pair of quarks, namely three potential energy terms. Hence we expect that baryons are more compact and deeply bound than diquarks both because of heavier reduced mass and due to more potential energy, as can be seen in Fig.4(left). The binding energy is slightly greater than temperature at T_c and comparable up to 1.3 T_c, which means the baryons (and anti-baryons) should play some role for temperatures not too high above T_c. As far as we know, this point is noticed and demonstrated for the first time.

Having studied baryons, we go a step further and include the simplest closed 3-chain structure above T_c, namely ggg, a color-singlet channel in which the three gluons mutually interact in similar way as three quarks in a baryon. In a string picture, there is a string between each pair, so the potential should be the same as in a meson (twice that for diquarks). Again we use the 3-body trial wavefunction (17) and minimize energy according to A, B, C. The results are shown in Fig.4(right). It is bound up to rather high

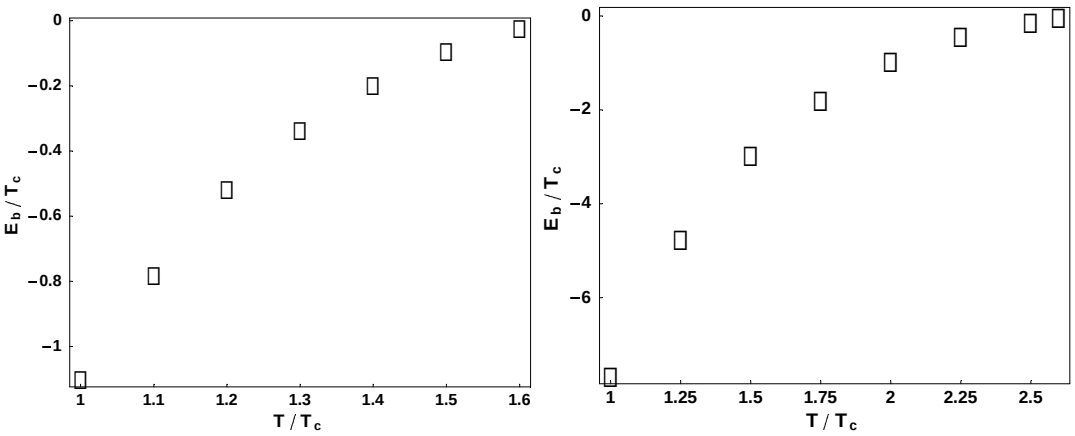

Figure 4. Dependence of the baryon (left) and trigluon (right) binding energy on the temperature. The units are in T_c.

temperature of about 2.6 T_c. The binding energy at T_c is as high as 7.64 T_c, and the size shrinks to only about 0.3 fm. Since we only use static potentials, without any relativistic corrections, we warn the reader that close to T_c the ggg binding (the only one!) becomes too large to be reliably evaluated inside the approximations made[3].

5. Summary and discussion

To sum up, three *multibody* bound states have been studied via variational approach: (i) "polymer chains" of the type $\bar{q}gg..gq$; (ii) baryons (qqq); (iii) closed (3-)chains of gluons (ggg). For the chains (i) we have proved that they have the same binding energy per bond as mesonic states and thus form in the same temperature range as mesonic states. We have established the binding energies and survival T-ranges for all these three structures. All the results are summarized in Fig.5 and Table.1. We conclude that between temperature region 1-1.5 T_c the existence of all these multibody bound states is not only possible but very robust.

Before we go forward with a general discussion, let us try to summarize the proposed scenario as a single picture, see Fig.6. From relatively short string-like configuration of color fields at low T, fig (a), one moves to longer strings (b) at the critical point [8]. New is picture (c) which depicts "polymeric chains" considered in this work, significant at $T = (1 - 1.5)T_c$. Eventually, at high T, one goes into (d) with independent quark and gluon quasiparticles, neutralized by isotropic Debye clouds.

We have not studied in this work neither more complicated states, such as a hybrid of baryons and polymers or a network of chains connected by color junctions, nor have we attempted to evaluate the possible role of polymers/baryons in thermodynamical and transport properties of sQGP (to be done elsewhere).

(In particular, we show in [16] that baryonic susceptibilities – up to the 6-th derivatives

[3]On top of relativistic effects we ignored, the simplification of the potential $log(1/r + 3T) \to log(1 + 3T)$ used above also affects binding close to T_c.

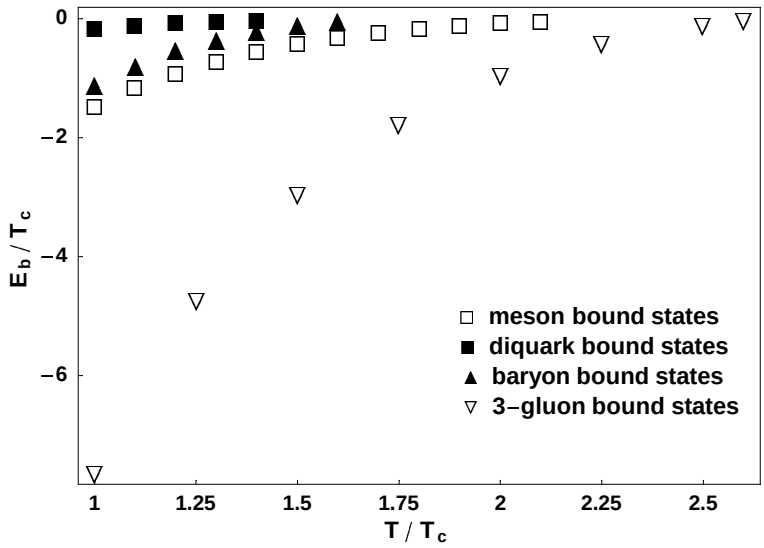

Figure 5. Dependence of various states' binding energy on the temperature. The units are in T_c.

Table 1
Summary of different bound states at $T > T_c$ studied in this paper. The column $C/C_{\bar{q}q}$ gives the relative potential strengths used in calculation, E_b means the binding energy, and T_m refers to the melting temperatures for different structures.

structure	-body	$C/C_{\bar{q}q}$	E_b/T_c at T_c	T_m
$\bar{q}q$	2	1	-1.45	2.1
$\bar{q}g\cdots gq$ (polymer chain)	N	1	-1.45∗(N-1)	2.1
ggg (closed chain)	3	1	-7.64	2.6
qq / $\bar{q}\bar{q}$	2	1/2	-0.13	1.4
qqq / $\bar{q}\bar{q}\bar{q}$ (closed chain)	3	1/2	-1.10	1.6

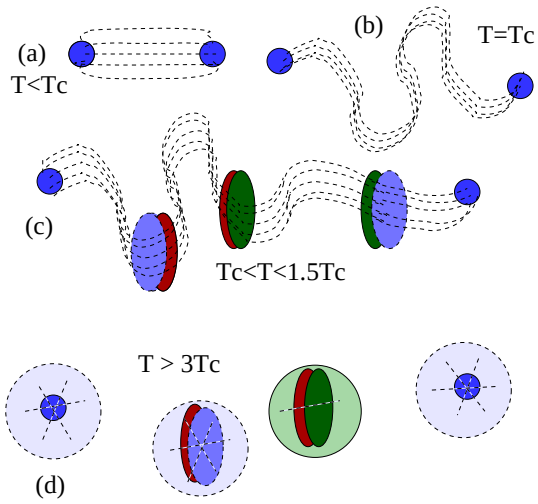

Figure 6. (color online) A schematic picture of the distribution of color fields, at different temperatures. Single color circles are quarks, bi-colored ovals are gluons.

over chemical potential – studied recently on the lattice by the UK-Bielefeld collaboration [12] provide good evidence in favor of existing of baryons near T_c, with a mass rapidly increasing into the QGP phase.)

Let us make here only few general comments on that. The conbtribution of polymers into partition function can be easily evaluated via a geometric series: the resulting enhancement factor, correcting the contribution of the $\bar{q}q$ mesonic states to those with any number of intermediate gluons $\bar{q}g...gq$, is

$$f_{polymers} = \frac{1}{1 - 6exp[(|\delta E| - M_g)/T]} \tag{18}$$

where 6 is the color and spin degeneracy added by each link. If one takes literally the bond binding we found, up to $1.4T_c$, and the effective gluon mass we used, one finds that even at $T = T_c$ this factor only reaches about 1.2, and rapidly decreases at higher T. The corresponding enhancement factor for baryons is a cube of that, but baryons themselves are a small effect. So, all this means that we think all multibody states we discussed provide just a few percent corrections into thermal properties of sQGP [4].

The reason we studied those states is to see if they can be more important for *transport properties*. It was first suggested by Shuryak and Zahed in [5] that existence of marginal states must increase rescattering and thus dramatically reduce viscosity (mean free paths), leading to a collisional (hydrodynamical) regime of expansion. Since different states get marginal at different T, one may hope this mechanism to work at all T up to about $2T_c$, the highest temperature at RHIC.

[4]However if some unincluded effects would increase this bidning by about a factor of two, the zero in the denominator can be reached, forcing total "polymerization" of matter.

We will now argue that the situation can be different for charm stopping or jet quenching. So far the estimates in the paper [13] were only done for jet energy losses due to ionization of *binary* mesonic states. The multibody bound clusters we studied above surely should induce jet quenching even more drastically. This should be especially true for *chain formation*, because it is known to be effective mechanism of the momentum distribution over larger volume of matter. The chains, connected by junctions, can in principle even produce a network.

The effect of similar substructure on transport is well known in material sciences: in fact an admixture of fibers is widely used to disperse forces applied to the system. The most famous examples are Kevlar fibers added to epoxy (or other plastic), now applied in wide range of applications, from tires, boats etc. to such exotic ones as "bullet-proof vests" and even "anti-mine boots".

Let us end up by speculating a bit about what would happen if indeed only the "well polymerized" sQGP in the interval $T = (1 - 1.5)T_c$ dominates the jet energy loss dE/dx. Since at RHIC the initial T is about $2T_c$, jets would have smaller losses till the matter cools down and polymerizes properly. Such a delay is observable because jets can move out of matter during such latent time: this would drastically affect the magnitude and especially angular distribution of the jet quenching.

In fact the original idea of "jet tomography" via jet quenching is in serious trouble for quite a while, because the most natural assumption – the energy loss proportional to the matter density – is in strong contradiction with the observed strong angular dependence of jet quenching [14], predicting too weak azimuthal asymmetry for non-central collisions. It was however recently pointed out by Pantuev [15] that a better description of data can be achieved if the jet quenching at the highest RHIC energy is switched on after some "latent time" of about 2.2 fm. This time quite reasonably matches a cooling time from $T \approx 2T_c$ to $T \approx 1.5T_c$ at RHIC.

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